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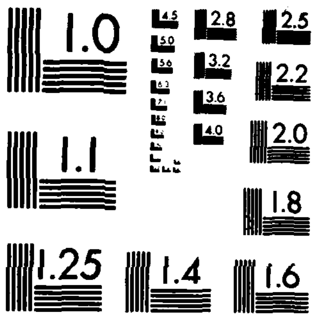

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We present an analytical expression for the growth rate of the $\underset{\sim}{E} \times \underset{\sim}{B}$ instability in the long wavelength limit, i.e., $k_y L_n < 1$, where k_y is the wavenumber and L_n is the scale length of the density gradient. Specifically, we obtain results for the case of a single discontinuity in plasma density.

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A

LONG WAVELENGTH LIMIT OF THE $\underset{\sim}{E} \times \underset{\sim}{B}$ INSTABILITY

I. Introduction

The $\underset{\sim}{E} \times \underset{\sim}{B}$ instability is regarded as an important instability in the structuring of ionospheric plasmas (e.g., the high latitude F region, barium clouds). The instability is basically an interchange mode and can be excited in a weakly collisional, inhomogeneous, magnetized plasma containing a neutral wind or an ambient electric field orthogonal to the magnetic field. Depending upon the ratios ν_e/Ω_e and ν_i/Ω_i, where $\nu_{e(i)}$ is the electron (ion) cyclotron frequency, two types of currents can be generated. For the case of $\nu_e/\Omega_e << 1$ and $\nu_i/\Omega_i << 1$, a Pedersen current is produced by the ions; when $\nu_e/\Omega_e << 1$ and $\nu_i/\Omega_i >> 1$, a Hall current is produced by the electrons. The Pedersen current driven instability (Simon, 1963; Hoh, 1963) is relevant to F region irregularities (Ossakow, 1979), while the Hall current driven instability (Rogister and D'Angelo, 1970; Sudan et al., 1973) is relevant to E region irregularities (Farley, 1979; Fejer and Kelley, 1980). This brief report will discuss the former instability.

An extensive amount of research has been devoted to the E x B instability, both theoretical (Linson and Workman, 1970; Shiau and Simon, 1972; Perkins et al., 1973; Ossakow et al., 1978; Huba et al., 1983) and numerical (Zabusky et al., 1973; McDonald et al., 1980; McDonald et al., 1981; Ossakow et al., 1982). In general, the geometry and plasma configuration used in the analyses are shown in Fig. 1a. The ambient magnetic and electric fields are in the z and y directions, respectively (i.e., $\underset{\sim}{B} = B_0\hat{e}_z$ and $\underset{\sim}{E} = + E_0\hat{e}_y$), and the density is inhomogeneous in the x direction [i.e., $n = n_0(x)$]. Wave perturbations are assumed to be primarily in the y direction so that $\delta p \propto \exp(ik_y y)$ where δp is some perturbed quantity. It is usually assumed that $k_y L_n >> 1$, where $L_n = (\partial \ell n\, n/\partial x)^{-1}$ is the scale length of the density gradient, and a <u>local</u> stability analysis is performed. [Perkins and Doles (1975) and Huba et al. (1983) are exceptions. They considered the effect of velocity shear on the $\underset{\sim}{E} \times \underset{\sim}{B}$ instability which required a <u>nonlocal</u> stability analyses]. The purpose of this brief report is to consider the opposite limit, viz., $k_y L_n << 1$, and to present an analytical expression for the growth rate.

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II. Theory

We use the plasma configuration and geometry shown in Fig. 1a and assume $\nu_e \ll \Omega_e$ and $\nu_i \ll \Omega_i$ so that an ion Pedersen drift exists in the y direction. We take perturbation quantities of the form $\delta p = \delta p(x) \exp[i(k_y y - \omega t)]$ and assume $k_y L_n \ll 1$ where $L_n = (\partial \ell n\ n/\partial x)^{-1}$ is the scale length of the density gradient, i.e., we are considering a discontinuous boundary layer.

The fundamental equations used in the analysis are continuity and momentum transfer:

$$\frac{\partial n_\alpha}{\partial t} + \nabla\cdot(n_\alpha \underset{\sim}{V}_\alpha) = 0 \tag{1}$$

$$0 = -\frac{e}{m_e}\left(\underset{\sim}{E} + \frac{1}{c}\underset{\sim}{V}_e \times \underset{\sim}{B}\right) \tag{2}$$

$$\frac{\partial \underset{\sim}{V}_i}{\partial t} = \frac{e}{m_i}\left(\underset{\sim}{E} + \frac{1}{c}\underset{\sim}{V}_i \times \underset{\sim}{B}\right) - \nu_{in}\underset{\sim}{V}_i \tag{3}$$

where α denotes species (e: electrons; i: ions) and other variables have their usual measuring. We neglect electron inertia but retain ion inertia. The equilibrium drifts are given by

$$\underset{\sim}{V}_e = 0 \tag{4}$$

$$\underset{\sim}{V}_i = \frac{\nu_{in}}{\Omega_i}\frac{cE_0}{B}\hat{e}_y \tag{5}$$

where we have chosen a reference frame such that $V_x = V_x - cE_0/B$ and $\Omega_i = eB/m_i c$.

We now consider a linear perturbation analysis of Eqs. (1)-(3). We assume $n_\alpha = n_\alpha + \delta n_\alpha$, $\underset{\sim}{V}_\alpha = \underset{\sim}{V}_\alpha + \delta\underset{\sim}{V}_\alpha$ and $\underset{\sim}{E} = \underset{\sim}{E}_0 - \nabla\phi$ where ϕ is the perturbed electrostatic potential. Using Eqs. (2) and (3) we find that

$$\delta\underset{\sim}{V}_e = -ik_y\phi\left(\frac{c}{B}\right)\hat{e}_x + \phi'\left(\frac{c}{B}\right)\hat{e}_y \tag{6}$$

and

$$\delta\underset{\sim}{V}_i = \frac{c}{B}\left[-ik_y\phi + i(\tilde{\omega}/\Omega_i)\phi'\right]\hat{e}_x$$

$$+ \frac{c}{B} \left[-k_y(\tilde{\omega}/\Omega_i)\phi + \phi'\right]\hat{e}_y \qquad (7)$$

where $\tilde{\omega} = \omega + i\nu_{in}$ and $\phi' = \partial\phi/\partial x$. Substituting Eqs. (6) and (7) into Eq. (1) and making use of quasi-neutrality ($\delta n_e = \delta n_i$) we obtain [Huba et al., 1983]

$$\phi'' + \frac{n'}{n}\phi' - \left[k_y^2 + \frac{k_y(cE_0/B)}{\tilde{\omega}} \frac{\nu_{in}}{\omega} k_y \frac{n'}{n}\right]\phi = 0 \qquad (8)$$

which we rewrite as

$$(n\phi')' - \left[nk_y^2 + \frac{k_y(cE_0/B)}{\tilde{\omega}} \frac{\nu_{in}}{\omega} k_y n'\right]\phi = 0 \qquad (9)$$

after multiplying through by n.

We now assume that

$$n = \begin{cases} n_1 & ; \quad x > 0 \\ n_2 & ; \quad x < 0 \end{cases} \qquad (10)$$

(as shown in Fig. 1b) since $k_y L_n \ll 1$, and take

$$\phi(x) = \phi_1 e^{-k_y x} + \phi_2 e^{k_y x} \qquad (11)$$

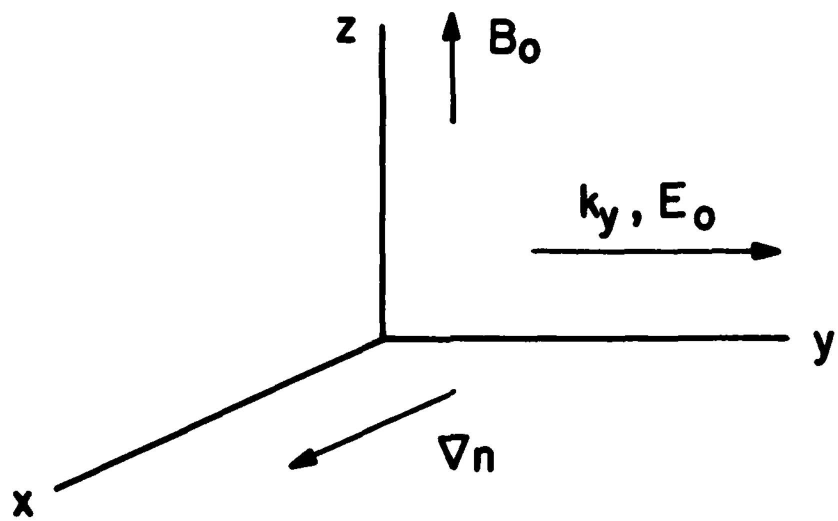
since Eq. (8) reduces to $\phi'' - k_y^2\phi = 0$ for $x \neq 0$. The modes are required to be bounded as $x \to \pm\infty$ so that

$$\phi(x) = \phi_1 e^{-k_y x} \qquad ; \quad x > 0 \qquad (12)$$

and

$$\phi(x) = \phi_2 e^{k_y x} \qquad ; \quad x < 0 \qquad (13)$$

We require that the interface velocity and the fluid velocity perpendicular to the interface be equal (Chandrasekhar, 1961) which requires that δV_x be continuous at the discontinuity, i.e., $x = 0$. From Eqs. (6) and (7) we find that this requires ϕ to be continuous at $x = 0$. Thus, $\phi_1 = \phi_2$ in Eqs. (12) and (13) so that



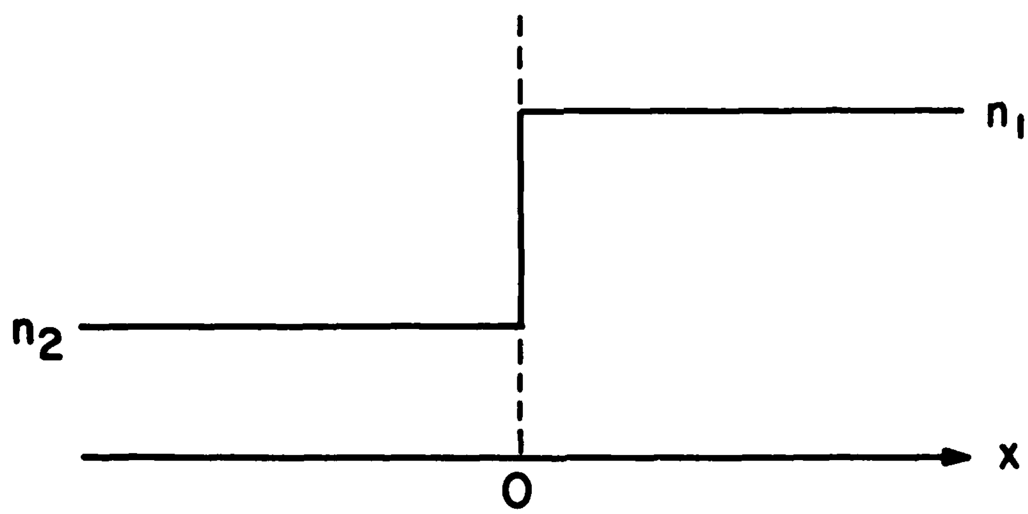


Fig. 1. Plasma geometry and slab configuration used in the analysis. (a) Standard plasma configuration. (b) Plasma configuration with a discontinuity in the density at x = 0.

$$\phi(x) = \begin{cases} \phi_0 \, e^{-k_y x} & ; \quad x > 0 \\ \phi_0 \, e^{k_y x} & ; \quad x < 0 \end{cases} \tag{14}$$

Finally, to obtain a dispersion equation for the modes we integrate Eq. (9) across the discontinuity at $x = 0$. Thus, we have

$$\int_{-\varepsilon}^{\varepsilon} (n\phi')' \, dx = \int_{-\varepsilon}^{\varepsilon} [nk_y^2 + \frac{k_y(cE_0/B)}{\tilde{\omega}} \frac{\nu_{in}}{\omega} k_y n']\phi \, dx \tag{15}$$

Since ϕ is continuous across the boundary at $x = 0$, it is found that Eq. (15) leads to

$$(n\phi')_1 - (n\phi')_2 = \frac{k_y(cE_0/B)}{\tilde{\omega}} \frac{\nu_{in}}{\omega} k_y(n_1\phi_1 - n_2\phi_2) \tag{16}$$

where (1, 2) indicate the region $x > 0$ (ε) and $x < 0$ ($-\varepsilon$), respectively. Substituting Eq. (14) into Eq. (16) and letting $\varepsilon \to 0$ we arrive at

$$\tilde{\omega} = -k_y(cE_0/B) \frac{\nu_{in}}{\omega} \frac{n_1 - n_2}{n_1 + n_2} \tag{17}$$

Equation (17) has the solution

$$\omega = -\frac{i\nu_{in}}{2} [1 \mp (1 + 4k_y \frac{(cE_0/B)}{\nu_{in}} \frac{n_1 - n_2}{n_1 + n_2})^{1/2}] \tag{18}$$

Equation (18) can be simplified by considering the limits $\omega \ll \nu_{in}$ and $\omega \gg \nu_{in}$, that is,

$$\omega = i \, k_y(cE_0/B) \frac{n_1 - n_2}{n_1 + n_2} \quad ; \quad \omega \ll \nu_{in} \tag{19}$$

and

$$\omega = i[k_y(cE_0/B)\nu_{in} \frac{n_1 - n_2}{n_1 + n_2}]^{1/2} \quad ; \quad \omega \gg \nu_{in} \tag{20}$$

so that instability results when $n_1 > n_2$ and $E_0 > 0$, i.e., $\omega = \omega_r + i\gamma$ with $\gamma > 0$.

We compare Eqs. (19) and (20) to the expressions obtained in the short wavelength limit ($k_y L_n \gg 1$). They are given by (Linson and Workman, 1970)

$$\omega = i\,\left(\frac{cE_0}{B}\right)\frac{1}{L_n} \qquad ; \quad \omega \ll \nu_{in} \qquad (21)$$

and (Ossakow et al., 1978)

$$\omega = i\left[\frac{cE_0}{B}\,\frac{\nu_{in}}{L_n}\right]^{1/2} \qquad ; \quad \omega \gg \nu_{in} \qquad (22)$$

Note that Eqs. (19) and (20) can be obtained from Eqs. (21) and (22) by making the identification

$$\frac{1}{L_n} \rightarrow k_y\,\frac{n_1 - n_2}{n_1 + n_2} \qquad (23)$$

in Eqs. (21) and (22). This identification (Eq. (23)) is the same one needed to make the transition from the short wavelength to long wavelength Rayleigh-Taylor instability (Chandrasekhar, 1961).

We note that we have not specified a form for the density perturbation δn. Also, we have made reference to the interface between the two materials without ever specifying either its initial perturbation or describing explicitly its subsequent evolution. Further, it is shown in Huba et al. (1983) that

$$\delta n_e = -\frac{c}{B}\,\frac{k_y\phi}{\omega}\,n' \qquad (24)$$

Since n' is a delta function at the interface, any finite amplitude perturbation such that ϕ were finite would cause δn_e to "blow up" at the interface, a nonsensical result.

The resolution of the above difficulties consists of starting with a sharp, but continuously differentiable, distribution of plasma of scale length L_n, complete with an infinitesimal perturbation of the form given by Eq. (24), and properly taking the limit of both the sharpness of the profile and the amplitude of the perturbation. Looking at Eq. (24), we see that δn_e is proportional to n'. That is to say, Eq. (24) is completely consistent with our describing the perturbation as a sinusoidal displacement of the fluid in the x direction. In fact, this displacement ξ is given by

$$\xi = \frac{c}{B}\frac{k_y\phi}{\omega} \tag{25}$$

The limiting process which avoids the problems noted above in defining δn_e for a discontinuity is to simultaneously let both L_n and ξ go to zero in a manner such that

$$\frac{\xi}{L_n} \rightarrow \text{constant as} \quad \begin{matrix} \xi \rightarrow 0 \\ L_n \rightarrow 0 \end{matrix} \tag{26}$$

Thus, the scale length over which the density changes from n_1 to n_2 and the displacement of the fluid are always of the same order, and δn_e remains bounded in the limit. The perturbation in this limit consists of displacing the interface sinusoidally, with the amplitude of this displacement growing in time with the growth rate given by Eq. (18).

III. Conclusion

We have presented an analytical expression for the growth rate of the $\underset{\sim}{E} \times \underset{\sim}{B}$ instability in the long wavelength limit, i.e., $k_y L_n << 1$, the limit of a single discontinuity in plasma density. The growth rate, in general, is given by Eq. (18). It is similar to that found in the short wavelength limit ($k_y L_n >> 1$) via the identification $L_n^{-1} \rightarrow k_y(n_1 - n_2)/(n_1 + n_2)$. We point out that Huba et al. (1983) recently investigated the $\underset{\sim}{E} \times \underset{\sim}{B}$ instability numerically. For a specific set of parameters [see Fig. 12 of Huba et al. (1983)] they find that $\gamma = 0.93\ k_y(cE_0/B)$ where $k_y L \simeq 0.1$. Using this same set of parameters in Eq. (18), we find that $\gamma = 0.95\ k_y(cE_0/B)$ which is in excellent agreement with the numerical results. We also note that the results presented here can be used to describe the bifurcation tendency of two-dimensional ionospheric barium clouds (McDonald et al., 1981; Overman and Zabusky, 1980) which will be reported in a future paper (Zalesak and Huba, 1983).

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